

Spread Supersymmetry

Lawrence J. Hall and Yasunori Nomura

*Berkeley Center for Theoretical Physics, Department of Physics,
and Theoretical Physics Group, Lawrence Berkeley National Laboratory,
University of California, Berkeley, CA 94720, USA*

Abstract

In the multiverse the scale of supersymmetry breaking, $\tilde{m} = F_X/M_*$, may scan and environmental constraints on the dark matter density may exclude a large range of \tilde{m} from the reheating temperature after inflation down to values that yield a lightest supersymmetric particle (LSP) mass of order a TeV. After selection effects, for example from the cosmological constant, the distribution for \tilde{m} in the region that gives a TeV LSP may prefer larger values. A single environmental constraint from dark matter can then lead to multi-component dark matter, including both axions and the LSP, giving a TeV-scale LSP somewhat lighter than the corresponding value for single-component LSP dark matter.

If supersymmetry breaking is mediated to the Standard Model sector at order $X^\dagger X$ and higher, only squarks, sleptons and one Higgs doublet acquire masses of order \tilde{m} . The gravitino mass is lighter by a factor of M_*/M_{Pl} and the gaugino masses are suppressed by a further loop factor. This Spread Supersymmetry spectrum has two versions, one with Higgsino masses arising from supergravity effects of order the gravitino mass giving a wino LSP, and another with the Higgsino masses generated radiatively from gaugino masses giving a Higgsino LSP. The environmental restriction on dark matter fixes the LSP mass to the TeV domain, so that the squark and slepton masses are order 10^3 TeV and 10^6 TeV in these two schemes. We study the spectrum, dark matter and collider signals of these two versions of Spread Supersymmetry. The Higgs boson is Standard Model-like and predicted to lie in the range 110 – 145 GeV; monochromatic photons in cosmic rays arise from dark matter annihilations in the halo; exotic short charged tracks occur at the LHC, at least for the wino LSP; and there are the eventual possibilities of direct detection of dark matter and detailed exploration of the TeV-scale states at a future linear collider. Gauge coupling unification is at least as precise as in minimal supersymmetric theories.

If supersymmetry breaking is also mediated at order X , a much less hierarchical spectrum results. The spectrum in this case is similar to that of the Minimal Supersymmetric Standard Model, but with the superpartner masses 1 – 2 orders of magnitude larger than those expected in natural theories.

1 Introduction

The physical origin of the weak scale is currently being probed at the Large Hadron Collider (LHC). In the Standard Model (SM), the weak scale is unnatural; in particular, the Higgs mass parameter must be fine-tuned by many orders of magnitude, leading to the almost universal belief that some new physics must be rigidly connected to the weak scale, $v \sim 200$ GeV. Supersymmetry, a rather unique and elegant extension of spacetime symmetry, has emerged as the leading candidate for this new physics, with the scale of weak interactions linked directly to the scale of supersymmetric particles, \tilde{m} .

However, the multiverse, suggested by the plethora of string theory vacua [1], casts doubt on this picture. The overall magnitude of the weak scale may be selected by anthropic requirements for an observable universe. Denying a symmetry explanation for the weak scale is controversial, but the exploration of an anthropic weak scale is well-motivated. While the physical effects that might restrict the size of the weak scale are not fully clear, we know that changing v by a factor of a few will lead to drastic changes of the universe [2]. Indeed, a similar argument applied to the cosmological constant has led to the successful understanding of the order of magnitude of the observed dark energy [3, 4]—something that has not been achieved using a symmetry argument.

Even if the weak scale is anthropically determined, the underlying theory may still be supersymmetric—after all, supersymmetry is necessary for the consistency of string theory. We are therefore faced with the interesting possibility that \tilde{m} and v are not rigidly coupled to each other, but rather are decoupled [5]. A fundamental Higgs boson with a mass very far below \tilde{m} would imply many orders of magnitude of fine-tuning which would be strong evidence for environmental selection of the weak scale. But with supersymmetry breaking out of reach of high energy colliders, what observations could confirm this picture?

Two schemes that allow precision tests of a finely tuned weak scale far below \tilde{m} are Split Supersymmetry [5] and High Scale Supersymmetry [6]. In Split Supersymmetry the scalar superpartners have masses near \tilde{m} while the fermionic superpartners are taken to have masses of order the TeV scale to account for the observed dark matter of the universe. Measurements of the gluino lifetime, the gauge Yukawa couplings of the Higgsinos, and the Higgs boson mass would all be correlated with \tilde{m} . Furthermore, such a correlation would imply that the Higgs is elementary up to \tilde{m} , so that the demonstration of fine-tuning would be convincing. Split Supersymmetry yields gauge coupling unification with precision comparable to the Minimal Supersymmetric Standard Model (MSSM).

In High Scale Supersymmetry all superpartners have masses of order \tilde{m} and are inaccessible to colliders. Nevertheless, for $\tilde{m} \gtrsim 10^{11}$ GeV the Higgs boson mass is predicted to be in the range of (128 – 141) GeV, depending on the composition of the Higgs boson. Moreover, many theories lead to values of the Higgs mass at the edge of this range, and the UV uncertainties from the

superpartner spectrum are extremely small, less than 0.5 GeV. In High Scale Supersymmetry it is the value of the Higgs boson mass itself that provides evidence of many orders of magnitude of fine-tuning. High Scale Supersymmetry can be made complete with axion dark matter; if f_a is above 10^{12} GeV, as typically expected, the axion misalignment angle is environmentally selected to be small [7, 8]. Gauge coupling unification, while less precise than in the MSSM, is nevertheless still significant.

In this paper we pursue two related studies. First we argue that in theories with \tilde{m} varying in the multiverse, an environmental requirement from Large Scale Structure forbids a very large window of \tilde{m} . This corresponds to the range in which the Lightest Observable-sector Supersymmetric Particle (LOSP) has a mass between order TeV and the reheating temperature of the universe after inflation T_R . For large values of T_R , this forbidden region divides theories into two very separate classes: those with (some) superpartners at the TeV scale and those without. Note that even in the former case, these superpartner masses are not directly linked to the scale of electroweak symmetry breaking and their existence is not needed to comprise dark matter. The Lightest Supersymmetric Particle (LSP), which may or may not be the LOSP, is cosmologically stable and contributes to dark matter; but typically we expect multi-component dark matter with a significant axion contribution. One member of this class has the entire MSSM in the TeV domain. This Environmental MSSM will typically have several orders of magnitude of fine-tuning in weak symmetry breaking, and likely accounts for only a fraction of dark matter.

Secondly, within the class of theories where the environmental constraint forces at least some superpartners to be at the TeV scale, we perform a top-down analysis to arrive at models with a very simple theoretical structure: Spread Supersymmetry. A key feature of this scheme is that the underlying structure of the theory forces a large spread in the superpartner spectrum so that, even though the LOSP (= LSP in this case) is in the TeV domain, \tilde{m} and most of the superpartner spectrum are several orders of magnitude larger than the TeV scale. Measurements on the few superpartners that have TeV scale masses correlate with the value of the Higgs boson mass, allowing \tilde{m} to be inferred and yielding evidence for a very high degree of fine-tuning.

The scheme of Spread Supersymmetry postulates that a chiral supermultiplet X responsible for supersymmetry breaking is charged under some symmetry, so that supersymmetry breaking is transferred to the MSSM sector via operators involving $X^\dagger X$, but not via operators linear in X , which would generate the gaugino masses, the supersymmetric Higgs mass (μ term), and the scalar trilinears (A terms). The leading supersymmetry breaking effects, therefore, arise from

$$\mathcal{L}_{\text{SB}} \sim \frac{1}{M_*^2} [X^\dagger X (Q^\dagger Q + U^\dagger U + D^\dagger D + L^\dagger L + E^\dagger E + H_u^\dagger H_u + H_d^\dagger H_d + H_u H_d)]_{\theta^4}, \quad (1)$$

where $\{Q, U, D, L, E\}$ and $\{H_u, H_d\}$ are the matter and Higgs superfields, and M_* is the scale at which supersymmetry breaking is mediated to the MSSM sector. The supersymmetry breaking masses generated by these operators are of order $\tilde{m} \equiv F_X/M_*$. Note that each operator has an

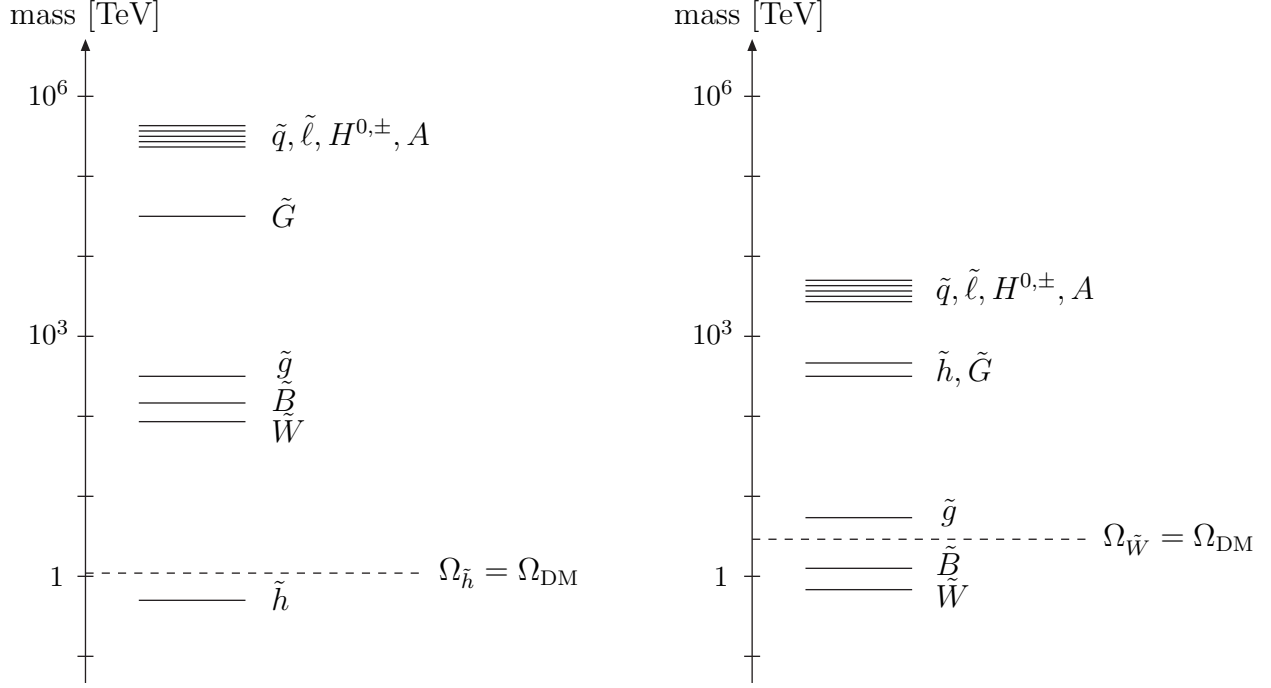


Figure 1: Two versions of the Spread Supersymmetry spectrum, with the Higgsino LSP (left) and the wino LSP (right).

unknown coefficient of order unity that is not displayed. The coefficient of the last term could be suppressed due to an approximate Peccei-Quinn (PQ) symmetry.

The gauginos and Higgsinos acquire masses from higher order effects in F_X/M_*^2 or through R symmetry breaking necessary to suppress the cosmological constant in supergravity.¹ This yields a spectrum for superpartners that spans some range. Since \tilde{m} scans in the multiverse, the forbidden window in the LOSP mass from the TeV scale to T_R will force either the LOSP to be heavier than T_R or to be in the TeV domain. The former simply gives a perturbation of High Scale Supersymmetry, but the latter leads to Spread Supersymmetry.

The spread spectrum we have in mind is illustrated in Fig. 1. The gravitino mass $m_{3/2} = F_X/\sqrt{3}M_{\text{Pl}} \equiv \epsilon_* \tilde{m}$ breaks R symmetry and anomaly mediation leads to gaugino masses of order $m_{3/2}/16\pi^2$, so that the gauginos are lighter than the squarks and sleptons by a factor $\epsilon_*/16\pi^2$. ($\epsilon_* \equiv M_*/\sqrt{3}M_{\text{Pl}}$ is typically smaller than 1.) The only remaining question is the mass of the Higgsinos, which is model dependent. Two versions of the spread spectrum are shown in Fig. 1. In the left panel the Higgsino masses arise from a one-loop radiative correction from virtual gauginos and Higgs bosons. In the right panel the Higgsino masses are of order the

¹We assume that the supersymmetry preserving vacuum expectation value $\langle X \rangle$ is sufficiently small that contributions to gaugino and Higgsino masses from operators involving $X^\dagger X$ are negligible.

gravitino mass, which can arise from supergravity interactions that follow from having $H_u H_d$ in the Kähler potential [9] or that cause a readjustment of the vacuum [10, 11]. The overall spread of the superpartner spectrum is typically 3 – 6 orders of magnitude. The normalization of the spectrum is set by environmental selection which forces the LSP mass below a critical value at the TeV scale. In both cases the spectra are sufficiently heavy to solve the supersymmetric flavor, CP and cosmological gravitino problems. The spectra shown are for $\epsilon_* \sim 10^{-2}$, corresponding to a high messenger scale of supersymmetry breaking of $M_* \sim 10^{16}$ GeV. Since the normalization of the spectra is determined by the environmental selection of dark matter, as M_* (and so ϵ_*) is reduced the only effect is to raise the masses of the squarks, sleptons and heavy Higgs—the gravitino, Higgsino and gaugino masses are not affected. For definiteness we will consider $\epsilon_* \sim 10^{-2}$ in the rest of the paper.

In the next section, we explore in a fairly general setting how environmental selection can exclude a very large range of \tilde{m} . In section 3, we study Spread Supersymmetry with the LSP Higgsino spectrum of the left panel of Fig. 1. We elucidate the theoretical structure of the model and study the phenomenology of the Higgsino states at colliders and for dark matter. In section 4 we introduce and study the model of Spread Supersymmetry with the LSP wino spectrum of the right panel of Fig. 1. For both theories we pay attention to the Higgs boson mass prediction, and we also point out that in both cases the precision of gauge coupling unification is comparable to that of the MSSM. In section 5 we comment on the Environmental MSSM. In section 6 we consider the case in which \tilde{m} and μ scan independently. Finally, we conclude in section 7.

2 The Forbidden Window in \tilde{m}

We consider any supersymmetric theory where the ratio of superpartner masses is fixed, but the overall scale of supersymmetry breaking, $\tilde{m} = F_X/M_*$, scans. Furthermore, we assume that the LSP is cosmologically stable, and that the reheating temperature after inflation, T_R , does not scan. In the multiverse, the value of \tilde{m} is determined by the prior distribution and anthropic selection. We assume that there is an anthropically allowed region for the amount of dark matter abundance; candidate boundaries for the region are discussed in Ref. [8] and used to limit the axion component of dark matter. This affects selection of \tilde{m} , since the LSP relic abundance depends on \tilde{m} .

The LSP may or may not be the LOSP—if not (e.g. if it is the gravitino \tilde{G}) then its abundance is determined by late decay of the LOSP: $\rho_{\text{LSP}} = (m_{\text{LSP}}/m_{\text{LOSP}})\rho_{\text{LOSP}}$. Here, ρ_i is the energy density of species i , and we assume that m_{LSP} and m_{LOSP} are not many orders of magnitude different. If the LOSP mass is below T_R , it is brought into thermal equilibrium and, as the temperature drops below its mass, freezes-out with an abundance $\xi_{\text{LOSP}} = \rho_{\text{LOSP}}/s$, where s is the total entropy. The resulting dark matter density is very high if the LOSP mass is well above

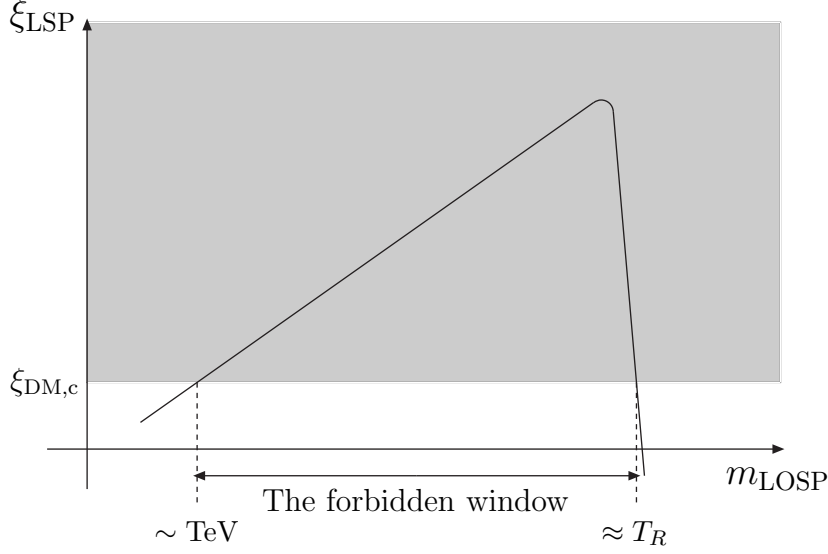


Figure 2: A large window between $\sim \text{TeV}$ and $\approx T_R$ for the LOSP mass, m_{LOSP} , is forbidden since it overproduces dark matter, beyond the bound in Eq. (2).

a TeV, beyond the upper edge of the anthropically allowed region characterized by a critical dark matter abundance $\xi_{\text{DM},c}$. This, therefore, leads to an environmental selection of \tilde{m}

$$\xi_a + \xi_{\text{LSP}}(\tilde{m}) < \xi_{\text{DM},c}, \quad (2)$$

where $\xi_{\text{LSP}} = (m_{\text{LSP}}/m_{\text{LOSP}})\xi_{\text{LOSP}}$, ξ_a is the axion abundance that depends on parameters in the axion sector, and we assume that possible dependence of $\xi_{\text{DM},c}$ on \tilde{m} is weak. This single condition simultaneously selects for a low enough axion density [7] and for a low LSP/LOSP mass. Hence, there is an environmentally forbidden window for \tilde{m} between the TeV scale and T_R , as illustrated in Fig. 2—the LOSP must either be heavy enough so that it is not produced significantly after inflation, or it must be light enough to satisfy Eq. (2).

The number of decades of this forbidden window increases with T_R as $\log_{10}(T_R/\text{TeV})$. For values of $T_R \gg \text{TeV}$, this window is very large and divides theories into two categories, those with no superpartners below T_R and those with at least some superpartners in the TeV domain. Many simple theories in the first category will yield High Scale Supersymmetry, with a Higgs mass prediction in the range of (128 – 141) GeV, depending on $\tan \beta$, providing T_R is 10^{10} GeV or larger. For theories in the second category, Large Scale Structure may not limit how light the LOSP can be, since dark matter may be fully accounted for by axions. Since supersymmetry has not yet been discovered at colliders, this suggests that the multiverse distribution for \tilde{m} in this region favors larger values, so that the LOSP mass is near the edge of the forbidden window. This second category of theories we call TeV-LOSP Supersymmetry.

The probability distribution for observing a universe with supersymmetry breaking \tilde{m} may then be written as $f(\tilde{m}) \theta(\tilde{m}) d\tilde{m}$, where θ is unity (vanishing) for values of \tilde{m} outside (inside) the forbidden window, and f contains effects from the a priori distribution of \tilde{m} in the multiverse and environmental selection other than Eq. (2).² If f increases with \tilde{m} for LOSPs in the TeV region, does this mean that High Scale Supersymmetry is more likely than TeV-LOSP Supersymmetry? No. Even if f increases very rapidly at low \tilde{m} , it is possible that f transitions to a decreasing (or less rapidly growing) function of \tilde{m} somewhere in the window because the forbidden region is very large. In particular, since the two anthropically allowed regions are very distant from each other, the size of the anthropic factor in the two regions may differ by many orders of magnitude. For example, if our universe has TeV scale superpartners, increasing \tilde{m} by many orders of magnitude will lead to a large change in the size of the low energy gauge couplings. Normalizing particle physics relative to the size of the QCD scale, this leads to a very large change in the strength of the gravitational interaction. The bottom line is that we do not know the relative probabilities of universes that are far apart in parameter space.

In Fig. 3 we illustrate how the superpartner spectrum depends on \tilde{m} for a spread spectrum (left panel), where supersymmetry breaking is transmitted to the superpartners at order $X^\dagger X$, and for a normal spectrum (right panel), where the transmission of supersymmetry breaking is linear in X . In the left panel we have chosen to show the spectrum corresponding to Spread Supersymmetry with a Higgsino LSP. For \tilde{m} well below the cutoff scale M_* , the spectra are simply proportional to \tilde{m} but, as \tilde{m} approaches M_* , higher dimension operators, giving the masses contributions of order $F_X^2/M_*^3 = \tilde{m}^2/M_*$, decrease the splitting of the spread spectrum and eventually the effective theory below M_* ceases to be supersymmetric.

The forbidden window is shown on each panel of Fig. 3 and corresponds to LOSP masses in the TeV to T_R range. Theories with \tilde{m} above the forbidden window are labeled as High Scale Supersymmetry in both panels. Providing T_R is large, the splitting of the superpartner spectrum provides only small corrections to the Higgs boson mass prediction. Theories with \tilde{m} below the forbidden window are labeled as Spread Supersymmetry in the left panel and Environmental MSSM in the right panel; they have very different experimental consequences. Spread Supersymmetry has the advantage that the squarks and leptons are sufficiently heavy to solve both the supersymmetric flavor and CP problems, and the gravitino is sufficiently heavy that there is no cosmological gravitino problem. Both these theories preserve successful gauge

²For $\tilde{m} \gg v$ the fine-tuning necessary for electroweak symmetry breaking contributes a factor \tilde{m}^{-2} to $f(\tilde{m})$. If this were the only effect, the a priori distribution must be peaked more strongly than \tilde{m}^2 in order for $f(\tilde{m})$ to prefer larger \tilde{m} values. It is, however, likely that other effects also contribute [12]. For example, if the cosmological constant scans and is selected by the environmental condition of Ref. [3], it strongly favors a large dark matter density—for single component LSP dark matter it provides a factor \tilde{m}^8 to $f(\tilde{m})$ —and it is likely to overwhelm both the a priori distribution and the effect from Higgs fine-tuning, yielding TeV-LOSP supersymmetry with dark matter close to the environmental boundary.

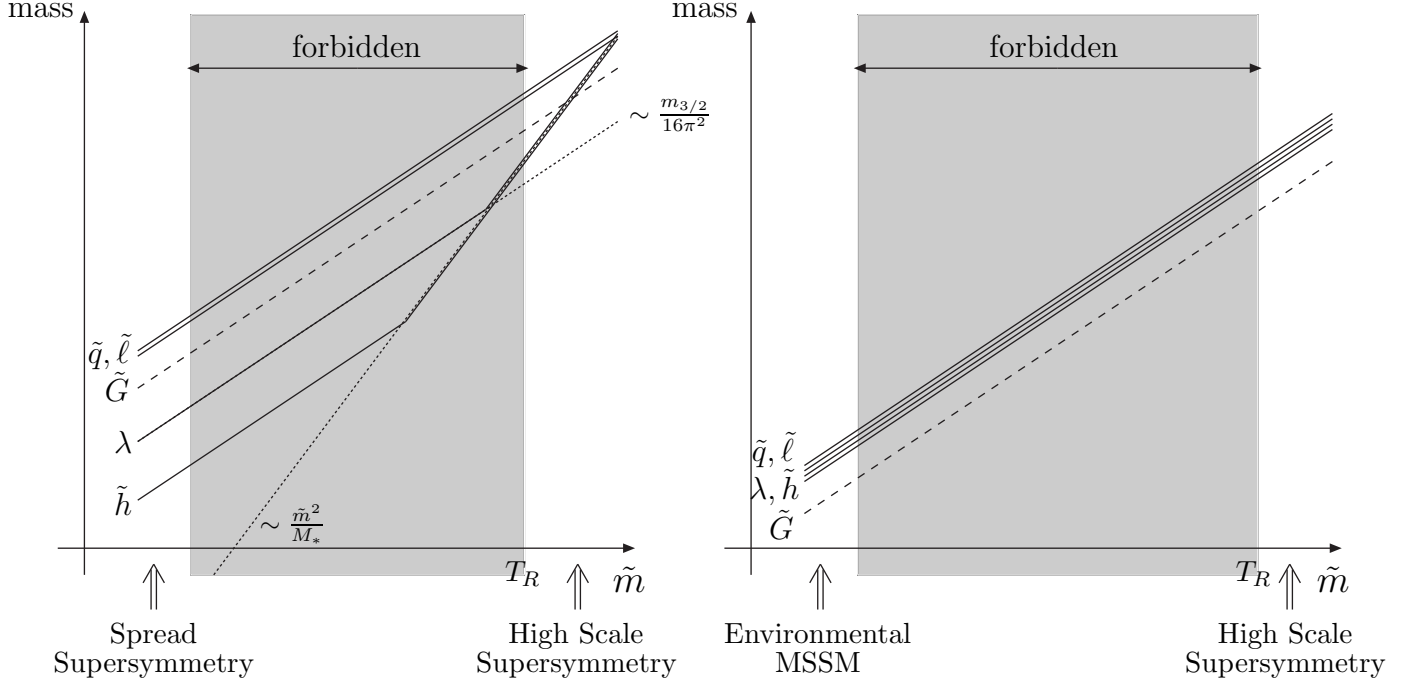


Figure 3: The masses of superpartners as a function of \tilde{m} , both on logarithmic scales, for the case where supersymmetry breaking is transmitted to the superpartners at order $X^\dagger X$ (left) and at order X (right). The forbidden window of Section 2 is indicated by the shaded areas.

coupling unification of the MSSM. We devote the next two sections to the Higgsino LSP and Wino LSP versions of Spread Supersymmetry, and we comment on Environmental MSSM in Section 5.

3 Spread Supersymmetry with Higgsino LSP

3.1 Spectrum

In any scheme with supersymmetry breaking arising at quadratic order in X , as in Eq. (1), squarks, sleptons and the heavy Higgs doublet, H , have masses of order \tilde{m} . The gauginos obtain masses of order $\epsilon_* \tilde{m}/16\pi^2$ from anomaly mediation [13, 14]; more specifically, $M_i = b_i g_i^2 m_{3/2}/16\pi^2$, where $(b_1, b_2, b_3) = (33/5, 1, -3)$ are the beta-function coefficients, and we have taken the phase convention that $M_{1,2}$ are real and positive. The lightest gaugino, therefore, is the wino.

The crucial question is how the Higgsino mass arises. In this section we assume that the supersymmetric term $H_u H_d$ is absent both in the Kähler potential and in the superpotential.

Since PQ symmetry is broken by the last term of Eq. (1) (and that R symmetry is broken by the gaugino masses), the Higgsino mass arises at one loop from a diagram with virtual electroweak gauginos and Higgs bosons, giving

$$m_{\tilde{h}} = -\frac{\sin 2\beta}{32\pi^2} \left(3g^2 M_2 \ln \frac{M_H}{M_2} + g'^2 M_1 \ln \frac{M_H}{M_1} \right), \quad (3)$$

where the light Higgs doublet is defined by $h = \sin \beta H_u + \cos \beta H_d$, with β in the first quadrant, and M_H , the mass of H , is expected to be of order \tilde{m} . (The sign convention of $m_{\tilde{h}}$ is such that it agrees with that of μ in Ref. [15] in the supersymmetric case.) This loop-suppressed Higgsino mass, together with the environmental selection discussed in the last section, leads to the spread spectrum shown in the left panel of Fig. 1.

The degeneracy between the charged Higgsino and neutral Higgsinos is lifted by electromagnetic corrections and via mixing with the gauginos, to yield the mass eigenstates χ_1^+ and $\chi_{1,2}^0$ with masses

$$m_{\chi_1^+} = |m_{\tilde{h}}| + \Delta_{\text{EM}} + \sin 2\beta \frac{M_W^2}{M_2}, \quad (4)$$

$$m_{\chi_{1,2}^0} = |m_{\tilde{h}}| \mp \left(\frac{1 \pm \sin 2\beta}{2} \right) \xi \frac{M_W^2}{M_2}, \quad (5)$$

where $\xi \equiv 1 + \tan^2 \theta_W (M_2/M_1)$ is numerically close to unity, and θ_W is the weak mixing angle. $\Delta_{\text{EM}} \simeq 310$ MeV represents the electromagnetic corrections [16, 17]. The negative sign in Eq. (3) is crucial, since it means that the charged Higgsino-wino mixing raises the mass of χ_1^+ so that the LSP is always neutral. The numerical size of the shifts of the Higgsino masses due to mixing with the gauginos is governed by the Higgsino mass via Eq. (3)

$$\frac{M_W^2}{M_2} \sim 300 \sin 2\beta \left(\frac{\text{TeV}}{|m_{\tilde{h}}|} \right) \text{ MeV}, \quad (6)$$

and, for $\sin 2\beta \sim 1$ and $|m_{\tilde{h}}| \sim \text{TeV}$, is the same order as the electromagnetic shift Δ_{EM} . The ordering of the spectrum is $m_{\chi_1^+} > m_{\chi_2^0} > m_{\chi_1^0}$. The three gaugino masses and three Higgsino masses depend on just two free parameters, $m_{3/2}$ and $\tan \beta$. (There is also a weak logarithmic dependence of the Higgsino masses on M_H .)

The present theory allows for a rather robust prediction of the Higgs boson mass as a function of angle β . This is because the sensitivity of the Higgs mass to the superpartner masses is fairly weak, as found in Ref. [6], and because the top-squark mixing parameter $\theta_{\tilde{t}} \approx A_t/m_{\tilde{t}}$ is very small, of order $\epsilon_*/16\pi^2$, as the only source of A terms is anomaly mediation, which is suppressed compared with the scalar masses. In Fig. 4, we show the Higgs boson mass as a function of $\sin 2\beta$, with the masses of the Higgsino, gaugino, and scalars taken to be 1 TeV, 10^5 GeV, and $\tilde{m} = 10^8$ GeV, respectively. The uncertainty coming from changing \tilde{m} by an order of magnitude in both directions is shown as the light-shaded band. The uncertainty from the gaugino masses is

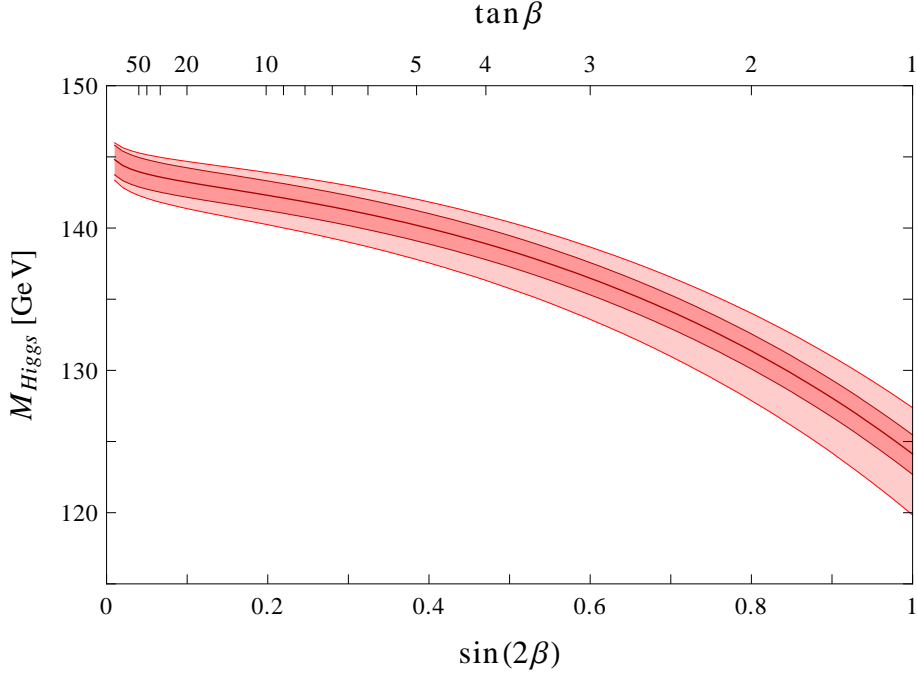


Figure 4: The Higgs mass prediction in Spread Supersymmetry with Higgsino LSP as a function of $\sin 2\beta$. (The corresponding values of $\tan \beta$ are also indicated at the top.) The solid red curve gives the Higgs mass prediction for $m_t = 173.2$ GeV, while the dark-shaded band shows the uncertainty coming from the experimental error of $\delta m_t = \pm 0.9$ GeV. The uncertainty from the superpartner mass scale \tilde{m} is depicted by the light-shaded band, which corresponds to $10^7 \text{ GeV} < \tilde{m} < 10^9 \text{ GeV}$.

of similar size, but generically smaller, and the effect of changing the Higgsino mass is negligible. For the top quark mass and QCD coupling, we have used the central values of the latest results: $m_t = 173.2 \pm 0.9$ GeV [18] and $\alpha_s(M_Z) = 0.1184 \pm 0.0007$ [19]. The uncertainty from the error of m_t is depicted by the dark-shaded band, with the width corresponding to the 1σ range. The uncertainty from α_s is smaller.

Note that the top Yukawa coupling y_t is asymptotically free in the SM, so its value at \tilde{m} is smaller than the low energy value. This allows for $\tan \beta = 1$, without encountering a Landau pole below the unification scale. In Fig. 5, we have plotted the running gauge (solid, blue) and top Yukawa (dashed, red) couplings as a function of energy for $\tan \beta = 1$. The jump in y_t comes from matching the one Higgs (below \tilde{m}) to two Higgs (above \tilde{m}) theories. The top Yukawa coupling is well perturbative at the unification scale. Requiring that all the Yukawa couplings are perturbative up to the unification scale, we find $0.7 \lesssim \tan \beta \lesssim 100$. Therefore, the range of $\sin 2\beta$ consistent with perturbative gauge coupling unification is

$$0.02 \lesssim \sin 2\beta \leq 1, \quad (7)$$

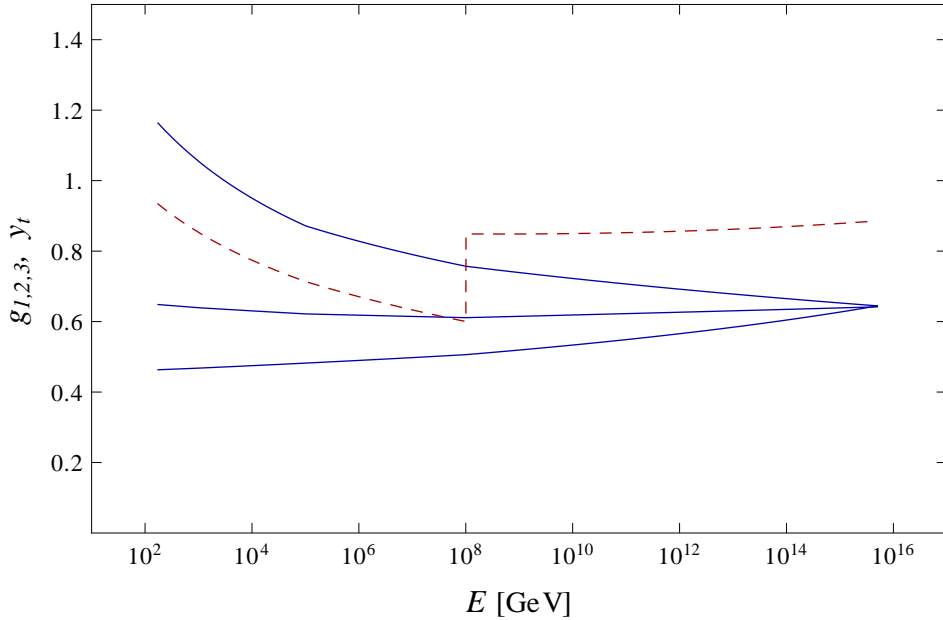


Figure 5: Evolution of the three gauge couplings, $g_{1,2,3}$ (solid, blue) and the top Yukawa coupling (dashed, red) for $\tan \beta = 1$ in Spread Supersymmetry with Higgsino LSP. The hypercharge gauge coupling has $SU(5)$ normalization.

which spans almost the entire range of Fig. 4.

3.2 Unification

As can be seen in Fig. 5, unification of the three SM gauge couplings works very well in Spread Supersymmetry. To quantify it, let us consider the size of the threshold correction $\delta(E)$ required for gauge coupling unification at energy E , where $\delta \equiv \sqrt{(g_1^2 - \bar{g}^2)^2 + (g_2^2 - \bar{g}^2)^2 + (g_3^2 - \bar{g}^2)^2} / \bar{g}^2$ with $\bar{g}^2 \equiv (g_1^2 + g_2^2 + g_3^2)/3$. In Spread Supersymmetry with Higgsino LSP, this quantity takes a minimum value at

$$M_{\text{unif}} \simeq (5 - 8) \times 10^{15} \text{ GeV}, \quad (8)$$

with

$$\delta(M_{\text{unif}}) \simeq 0.004 - 0.008, \quad (9)$$

where the value of M_{unif} is most sensitive to the gaugino masses (anti-correlation), while that of $\delta(M_{\text{unif}})$ to the Higgsino mass (correlation). The result of Eq. (9) can be compared with the values in the SM and in the MSSM: $\delta_{\text{min, SM}} \simeq 0.06$ and $\delta_{\text{min, MSSM}} \simeq O(0.01)$ (depending on the superpartner spectrum). Spread Supersymmetry, therefore, achieves gauge coupling unification, at least, at a level of the MSSM.

The value of the unified gauge coupling at $E \simeq M_{\text{unif}}$ is

$$g_{\text{unif}}(M_{\text{unif}}) \simeq 0.65. \quad (10)$$

This and Eq. (8) have an important implication on the rate of dimension six proton decay, caused by an exchange of the unified gauge bosons. Since the (partial) decay rate is proportional to $g_{\text{unif}}(M_{\text{unif}})^4/M_{\text{unif}}^4$, we find

$$\frac{\Gamma_{\text{Spread}}}{\Gamma_{\text{MSSM}}} \simeq 30 - 200, \quad (11)$$

where we have used $g_{\text{unif}}(M_{\text{unif}})|_{\text{MSSM}} \simeq 0.7$ and $M_{\text{unif}}|_{\text{MSSM}} \simeq 2 \times 10^{16}$ GeV. This corresponds to the lifetime [20]

$$\tau_{p \rightarrow e^+ \pi^0} \simeq (0.8 - 5) \times 10^{34} \text{ years} \quad (12)$$

in 4-dimensional supersymmetric grand unified theories. This range is just above the current lower limit from Super-Kamiokande $\tau_{p \rightarrow e^+ \pi^0} > 8.2 \times 10^{33}$ years [21], and can be fully covered by the planned Hyper-Kamiokande experiment at the 3σ level [22]. If grand unification is realized in higher dimensions [23], dimension six proton decay can have a variety of final states [24]. The lifetime in this case is also expected to be shorter than the corresponding case in which the low energy theory is the MSSM.

3.3 Dark matter

If dark matter is composed only of Higgsinos then the freeze-out mechanism requires the Higgsino mass to be 1.1 TeV. On the other hand, with multi-component dark matter the Higgsino fraction, and therefore the Higgsino mass, depends on the relevant multiverse distribution functions. This makes the mass of the Higgsino lighter:

$$m_{\tilde{h}} \simeq 1.1 \left(\frac{\Omega_{\tilde{h}}}{\Omega_{\text{DM}}} \right)^{1/2} \text{ TeV}. \quad (13)$$

(In this subsection, we ignore the small difference between $m_{\tilde{h}}$ and the LSP mass $m_{\chi_1^0}$.)

For illustration, let us take a distribution function $f(\tilde{m})d\tilde{m} = \tilde{m}^p(d\tilde{m}/\tilde{m})$ for values of \tilde{m} giving a Higgsino mass in the region corresponding to the critical environmental boundary of Eq. (2). This includes a quadratic weighting factor resulting from the environmental requirement that the Higgs vacuum expectation value is below its critical value. For mixed Higgsino/axion dark matter, the environmental boundary of Eq. (2) can be cast in the form

$$x^2 + y^2 < 1; \quad x = \frac{\tilde{m}}{\tilde{m}_c} = \frac{m_{\tilde{h}}}{m_{\tilde{h}_c}}, \quad y = \frac{\theta}{\theta_c}, \quad (14)$$

where θ is the axion misalignment angle, whose multiverse distribution is expected to be flat, and subscripts refer to the critical values. As an example, $p = 1$ leads to equal multiverse

averages of Higgsino and axion dark matter: $\langle\Omega_{\tilde{h}}\rangle = \langle\Omega_a\rangle$, giving an average Higgsino mass of about 700 GeV. Lower values of p give lighter Higgsinos; if $p = \epsilon \ll 1$, Higgsinos account on average for only a fraction ϵ of the dark matter, giving an expected value of the Higgsino mass of $\approx \sqrt{\epsilon}$ TeV. Ultimately, experiments will see a correlation between the Higgsino mass and the fractions of Higgsino and axion dark matter.

In Spread Supersymmetry with a Higgsino-like LSP, the direct detection of dark matter is challenging. The mass splitting between the various components of the Higgsino, Eqs. (4) and (5), are sufficient that the LSP-nucleus scattering occurs as the elastic scatter of a Majorana fermion. In the limit that the gaugino component of χ_1^0 is ignored, the spin-independent cross section for scattering from nuclei has been studied in Refs. [16, 25, 26]. In Ref. [16], 1-loop electroweak gauge boson contributions (from both box and Higgs exchange diagrams) led to a scattering cross section from protons of $\sigma_p \sim 2 \times 10^{-46} \text{ cm}^2$ for a Higgsino mass of 1 TeV and a Higgs mass of 115 GeV. In Ref. [25], diagrams involving gluons were included leading to a similar result for σ_p . However, the authors of Ref. [26] find an additional twist 2 operator contribution, and disagree with certain results of the earlier papers, concluding that for Higgs masses in the predicted range of Fig. 4 there is a cancellation between the Higgs exchange and non-Higgs exchange diagrams, leading to $\sigma_p \sim 10^{-48} \text{ cm}^2$.

Our LSP, however, contains a gaugino component, so additional contributions to the spin-independent scattering from nuclei arise from a tree-level Higgs exchange amplitude proportional to the small wino component. This contribution depends on $\tan\beta$ and has been studied in Ref. [27]. For gaugino masses of interest to us, this contribution to σ_p is of order 10^{-47} cm^2 or smaller. Further study should investigate the effect of uncertainties of the nuclear matrix elements on cancellations between the varying amplitudes. The spin-dependent cross section is generically dominated at tree level, and is in the range of $O(10^{-46} - 10^{-44} \text{ cm}^2)$ [27].

The indirect detection of Higgsino dark matter may occur via the detection of monochromatic photons of energy $m_{\tilde{h}}/2$ arising from $\tilde{h}\tilde{h} \rightarrow \gamma\gamma$ in the halo of our galaxy. For $m_{\tilde{h}} = 100 \text{ GeV} - 1 \text{ TeV}$, the one-loop electroweak annihilation cross section is $\sigma v(\tilde{h}\tilde{h} \rightarrow \gamma\gamma) \approx 10^{-28} \text{ cm}^3/\text{s}$ and has no significant Sommerfeld enhancement. Using the NFW profile for the distribution of galactic dark matter, data from the Fermi LAT place an upper limit on this cross section of $(20 - 100) \times 10^{-28} \text{ cm}^3/\text{s}$ for $m_{\tilde{h}}$ in the range of $(100 - 300) \text{ GeV}$ [28]. The eventual discovery of a monochromatic galactic photon signal would directly yield the Higgsino mass, motivating the construction of a linear lepton collider with a definite energy. In addition, from this mass one could accurately infer the fraction of dark matter in LSPs.

3.4 Collider signals of Higgsinos

Pairs of $(\chi_1^+, \chi_1^0, \chi_2^0)$ will be produced by the Drell-Yan mechanism at LHC. The charged state beta decays to the neutral states, but the mass difference is greater than ≈ 300 MeV allowing the two-body decay $\chi_1^+ \rightarrow \chi_{1,2}^0 \pi^+$. This gives $c\tau \lesssim 1$ cm, so that the charged tracks of χ_1^+ are not visible. For $|m_{\tilde{h}}| \sim \text{TeV}$, the pion is too soft to see above backgrounds. Similarly, the decays $\chi_2^0 \rightarrow \chi_1^0 e^+ e^-$ give $e^+ e^-$ pairs that are too soft to see above backgrounds. However, for lower $|m_{\tilde{h}}|$, the mass splittings between the states may be large for $\sin 2\beta \sim 1$ (see Eq. (6)), so that events with decays of boosted χ_1^+ or χ_2^0 states could give observable LHC signals.

At a future lepton collider, signals of nearly degenerate Higgsinos result from $\chi_1^+ \chi_1^- \gamma$ and $\chi_1^0 \chi_2^0 \gamma$ productions followed by $\chi_1^\pm \rightarrow \chi_{1,2}^0 \pi^\pm$ and $\chi_2^0 \rightarrow \chi_1^0 e^+ e^-$ decays [29]. Hard initial state photon radiation is required since otherwise the event contains only the soft products of χ_1^+ and χ_2^0 decays and such events are overwhelmed by underlying events from the collision of beamstrahlung photons that produce soft particles. The SM background arising from $\nu \bar{\nu} \gamma$ production has a cross section about three orders of magnitude larger than that of $\chi_1^+ \chi_1^- \gamma$ and $\chi_1^0 \chi_2^0 \gamma$ productions for a Higgsino mass in the region of 1 TeV. Although this background can be reduced by using polarized beams, search strategies should be devised that involve the soft π^+ and $e^+ e^-$ from χ_1^+ and χ_2^0 decays. Although χ_1^+ has $c\tau \sim \text{cm}$, it may be possible to observe the charged tracks if the boost factor $\sqrt{s}/2m_{\tilde{h}}$ is sufficiently large and if the luminosity is large enough to yield events with decays occurring at times of a few τ .

Precision measurements of the overall Higgsino mass and the two mass splittings, Eqs. (4) and (5), would be a powerful probe of the theory. These three observables depend on only $m_{3/2}$, $\sin 2\beta$ and $\ln M_H/M_2$, which would be fit to the data. Furthermore, $\sin 2\beta$ is correlated with the Higgs mass prediction, Fig. 4, yielding a consistency check. One would then infer the gaugino spectrum. Observing splittings amongst the Higgsinos of a few hundred MeV would imply electroweak gaugino masses of order 100 TeV independent of the underlying theory. Such a large breaking of supersymmetry would imply a very large, but model dependent, fine-tuning in electroweak symmetry breaking. Within Spread Supersymmetry the amount of fine-tuning depends on the value of ϵ_* inferred from $m_{3/2}$ and $\ln M_H/M_2$. For $\epsilon_* \sim 10^{-2}$ the tuning is of order 1 part in $\approx 10^{14}$.

4 Spread Supersymmetry with Wino LSP

In this section we consider the case where supergravity effects yield a Higgsino mass of order $m_{3/2}$. This can arise from a Higgs bilinear term in the Kähler potential $K = \lambda H_u H_d + \text{h.c.}$ [9] or by a vacuum readjustment induced by supergravity corrections to the potential of flat super-

symmetry [10, 11]. The superparticle spectrum becomes that in the right panel of Fig. 1:

$$m_{\tilde{q}, \tilde{l}, H^0, \pm, A} \sim \tilde{m}, \quad m_{\tilde{h}, \tilde{G}} \sim \epsilon_* \tilde{m}, \quad m_{\tilde{g}, \tilde{W}, \tilde{B}} \sim \frac{\epsilon_*}{16\pi^2} \tilde{m}, \quad (15)$$

where the gaugino masses are generated by anomaly mediation as well as one loop of the Higgs-Higgsino. This, therefore, leads to a scenario similar to the one discussed in Refs. [14, 30].³

We assume that the wino is the LSP, which is the case unless the contribution from a Higgs-Higgsino loop dominates over that from anomaly mediation. Phenomenology of this theory is mostly as discussed in Ref. [30]. The Higgs boson mass depends on the gaugino and Higgsino masses as well as $\tan \beta$ and ϵ_* , and is generically in the range

$$M_{\text{Higgs}} \approx (110 - 140) \text{ GeV}. \quad (16)$$

Gauge coupling unification works essentially as in the MSSM; we find

$$\delta(M_{\text{unif}}) \lesssim O(0.01), \quad M_{\text{unif}} \simeq 10^{16} \text{ GeV}. \quad (17)$$

The lightest gaugino is typically the wino. The mass splitting between the charged and neutral components, $\Delta m \equiv m_{\chi^+} - m_{\chi^0}$, is given by

$$\Delta m = \Delta_{\text{EM}} + O\left(\frac{m_W^2 m_{\tilde{W}}}{m_h^2}\right) \simeq \Delta_{\text{EM}}, \quad (18)$$

where $\Delta_{\text{EM}} \simeq 160 \text{ MeV}$ is the electromagnetic contribution. The second contribution to Δm in Eq. (18) from mixing with the Higgsino is expected to be less than 1 MeV in our framework. If the wino LSP composes the entire dark matter, then $m_{\tilde{W}} \simeq (2.7 - 3.0) \text{ TeV}$ [31]. If the axion composes a significant part of dark matter, the mass of the wino becomes correspondingly smaller:

$$m_{\tilde{W}} \simeq (2.7 - 3.0) \left(\frac{\Omega_{\tilde{W}}}{\Omega_{\text{DM}}}\right)^{1/2} \text{ TeV}, \quad (19)$$

ignoring the variation of the Sommerfeld enhancement factor with the LSP mass. The fraction $\Omega_{\tilde{W}}/\Omega_{\text{DM}}$ is determined by the multiverse distribution function $f(\tilde{m})$.

The cross section for the direct detection of wino dark matter is generically an order of magnitude larger than for Higgsino dark matter. In the limit of neglecting the Higgsino component of the LSP, and ignoring contributions from operators involving gluons, Ref. [16] quotes a spin-independent cross section from the proton of $\sigma_p \sim 10^{-45} \text{ cm}^2$ for a wino mass of 2.4 TeV and a Higgs mass of 115 GeV. However, including the gluon operators the authors of Refs. [32, 26] find a cancellation between the Higgs exchange and non-Higgs exchange diagrams, leading to

³These papers did not discriminate between the masses of scalars and the gravitino, $m_{\tilde{q}, \tilde{l}, H^0, \pm, A} \sim m_{\tilde{G}}$, corresponding to the case $\epsilon_* \approx O(1)$ in our scenario.

$\sigma_p \sim O(10^{-47} \text{ cm}^2)$ for a Higgs mass in the range of $\approx (114 - 140) \text{ GeV}$ and $m_{\tilde{W}} \sim 3 \text{ TeV}$. The tree-level Higgs exchange contribution, involving mixing with the Higgsino, is very small, $\lesssim 10^{-48} \text{ cm}^2$, in the parameter region of interest. The spin-dependent cross section is dominated by loop diagrams and is in the range of $O(10^{-46} - 10^{-44}) \text{ cm}^2$ [26].

The indirect detection of wino dark matter from annihilation to photons is promising. The cross section $\sigma v(\tilde{W}\tilde{W} \rightarrow \gamma\gamma)$ at one loop is about $10^{-27} \text{ cm}^3/\text{s}$, about an order of magnitude larger than for the case of pure Higgsino dark matter. Furthermore, while the Sommerfeld enhancement occurs at too large a mass to be relevant for Higgsinos, in the wino case the resonance occurs at a wino mass of about 2.3 TeV and yields an enhancement of the cross section by factors of (3, 30, 125) for wino masses of (1, 2, 2.5) TeV, respectively [33]. In both the Higgsino and wino LSP cases, a measurement of the energy of the photon line allows an inference of the fraction of dark matter carried by the LSP and motivates the construction of a lepton collider.

Prospects for detecting the wino LSP at the LHC are better than that for the Higgsino LSP. This is mainly because the mass splitting between the charged and neutral components is smaller, $\Delta m \simeq 160 \text{ MeV}$, so that the decay length of $\chi^+ \rightarrow \chi^0 \pi^+$ is longer, $c\tau \approx O(10 \text{ cm})$. Drell-Yan production of $\chi^+ \chi^-$ yields events with (disappearing) charged tracks, that can be triggered by high p_T jets or missing transverse energy [34, 17]. For $\sqrt{s} = 13 \text{ TeV}$ and integrated luminosity of 10 fb^{-1} (100 fb^{-1}), reach for the wino mass is estimated to be $\approx 350 \text{ GeV}$ (550 GeV) [17].

The wino could be discovered at a future lepton collider via very distinctive events arising from the production of $\chi^+ \chi^-$ followed by the decay $\chi^\pm \rightarrow \chi^0 \pi^\pm$ [35]. The highly ionizing χ^\pm tracks are crucial, since soft $\pi^+ \pi^-$ events have a large background from two photon induced processes. With $c\tau \sim 10 \text{ cm}$, independent of the wino mass, these ionizing tracks will be seen to end. Furthermore, a very soft pion will intersect the end of the track at large angle, and such events have large missing energy. Measuring the wino mass would allow the fraction of dark matter in winos to be inferred. Measuring the χ^+ lifetime would lead to a determination of Δm limiting the size of the Higgsino mixing contribution to the wino mass. For example, an observation that $\Delta m < 170 \text{ MeV}$ would imply that the Higgsino is heavier than the wino by at least a factor of $30\sqrt{m_{\tilde{W}}/\text{TeV}}$ and hence that there is a very high degree of fine-tuning in electroweak symmetry breaking, greater than one part in $10^5 (m_{\tilde{W}}/\text{TeV})^3$.

5 The Environmental MSSM

Suppose the supersymmetry breaking field X is neutral, rather than charged, so that all SM superpartners (including the Higgsino) acquire masses directly from the hidden sector at order \tilde{m} , through operators suppressed by M_* . In general, there can be mild hierarchies among these superpartners, so that there will be a great deal of model dependence in the precise superpartner

spectrum and the nature of the LOSP. It is, however, still true that there is generically a large forbidden window of \tilde{m} , as illustrated in the right panel of Fig. 3.

The key point is that the normalization of the spectrum will arise from the environmental bound resulting from the forbidden shaded zone of Fig. 3. The gravitino mass is $m_{3/2} = F_X/\sqrt{3}M_{\text{Pl}} \equiv \epsilon_*\tilde{m}$, and we expect the gravitino to be the LSP for all but the highest values of the messenger scale, $M_* \gtrsim \sqrt{3}M_{\text{Pl}}$. As ϵ_* is reduced, the gravitino mass gets smaller, implying that the environmental bound on the LOSP freeze-out abundance gets milder and the entire SM superpartners spectrum gets heavier. We do not consider this limit, since the bound on the reheating temperature from thermal production of gravitinos becomes more powerful $T_R < 10^9 \text{ GeV}(m_{3/2}/10 \text{ GeV})$, narrowing the range of the forbidden window. For $\epsilon_* \sim 10^{-2}$ the SM superpartner masses are all expected to be 1 to 2 orders of magnitude above the weak scale. Hence the lightest Higgs boson mass is predicted to be roughly in the range $\approx 100 - 130 \text{ GeV}$.

As one example of parameters for the Environmental MSSM, consider $\tilde{m} \sim 10 \text{ TeV}$, a LOSP mass somewhat larger than a TeV and $\epsilon_* \sim 10^{-3}$, giving $m_{3/2} \sim 10 \text{ GeV}$. For LOSP masses lighter than about a TeV, gravitinos arising from LOSP freeze-out and decay cannot comprise all the dark matter, because the LOSP decay upsets Big Bang Nucleosynthesis (BBN) [36]. In our scheme, however, the LOSP mass may be lighter than a TeV without being excluded by BBN because the gravitino may only be a sub-dominant component of dark matter. For LOSP masses above a TeV, dark matter could dominantly be gravitinos arising from LOSP freeze-out and decay, since the LOSP decays more rapidly and evades the BBN limits. Such decays of a $\tilde{\tau}$ LOSP could solve the BBN lithium problem [37].

If $\epsilon_* \gtrsim 1$, the LSP can be the LOSP—the superpartner spectrum is then as in the standard MSSM with gravity mediation (defined broadly) [38]. The scale of the superparticles, however, is now determined not by fine-tuning of electroweak symmetry breaking, but by the LSP relic abundance: $\Omega_{\text{LSP}} < \Omega_{\text{DM}}$. If the LSP contains a significant amount of the Higgsino or wino, this can lead to relatively heavy superparticles, which however may still be within reach of the LHC.

6 Scanning of μ

So far, we have assumed that the supersymmetric Higgsino mass parameter, μ , is absent from the superpotential. In this section we consider scanning μ and \tilde{m} independently.

If the multiverse distribution is logarithmic in μ (or favors small values), then typical observers live in universes with $\mu \ll \tilde{m}$ because electroweak symmetry does not break for $\mu \gtrsim \tilde{m}$ (so that no/few observers arise). In this case the Higgsino mass will be dominated by other sources. An effective μ term arises either from $K = \lambda H_u H_d$ or vacuum readjustment, leading to the theory discussed in Section 4, or is generated radiatively, leading to the theory described in section 3.

What if the a priori multiverse distribution for μ favors large μ ? In this case typical observers will see μ of order \tilde{m} , close to the upper limit imposed by the requirement of electroweak symmetry breaking. This is a multiverse solution to the μ problem and, in the context of Spread Supersymmetry (i.e. if X is charged), yields a wino LSP with a supersymmetric spectrum as illustrated in the right panel of Fig. 1, except that the Higgsinos now have a mass comparable to the scalar superpartners rather than to the gravitino. The collider and dark matter phenomenology is as in section 4. If X is neutral, the resulting theory is the Environmental MSSM discussed in the previous section. In either theory, the distribution for μ now contributes to the effective multiverse distribution that favors large values for supersymmetry breaking. The fine-tuning in the Higgs mass parameter can be overcome by a combination of the μ and \tilde{m} distributions, making it more plausible that \tilde{m} is typically (much) larger than the weak scale.

7 Conclusions

If a Standard Model Higgs is discovered at the LHC in the region of 115 – 145 GeV, with no signs of any other new physics, the possibility that the weak scale is determined by environmental selection on a multiverse will be increased. Does this mean that no new physics is expected in the TeV domain?

An environmental upper bound on the amount of dark matter in the universe can be phrased as an upper limit on the temperature of matter-radiation equality

$$T_{\text{eq}} < T_{\text{eq,c}}, \quad (20)$$

with the critical value, $T_{\text{eq,c}}$, not far above the observed value. In this case, any freeze-out relic of mass m , with annihilation cross section within a few orders of magnitude of m^{-2} , will satisfy $m < m_c \approx \text{TeV}$. Furthermore, if the multiverse distribution, including selection effects from other boundaries, favors large values of m , typical observers will find the mass of this weakly interacting massive particle (WIMP) close to its critical value. This quite generally leads to the possibility of new physics associated with WIMPs at the TeV scale, without having a direct connection to electroweak symmetry breaking. A key difference from conventional WIMPs is that the environmental boundary of Eq. (20) limits all cold relics suggesting multi-component dark matter, for example axions and WIMPs.

Specializing to supersymmetric theories, if the overall scale of supersymmetry breaking scans, with fixed superpartner mass ratios, the normalization of the spectrum may be determined by Eq. (20), fixing the LSP mass to be of order 1 TeV. This could lead to “Environmental Supersymmetry” with a spectrum of superpartners in the multi-TeV domain, 1 – 2 orders of magnitude larger than expected from natural theories of weak-scale supersymmetry. Alternatively, if the

leading supersymmetry breaking arises at quadratic order in the F_X spurion, a hierarchical spectrum of superpartners results, giving “Spread Supersymmetry.” We have studied the spectra, dark matter and collider signals of two examples of Spread Supersymmetry having Higgsino and wino LSPs with scalar superpartner masses of order 10^6 and 10^4 TeV, respectively. The examples studied have the same set of fields as the MSSM, but this is not necessary; a similar pattern of spectra would result in extended theories, for example with singlets added.

Signals for such theories may show up first via indirect detection of galactic dark matter through a monochromatic photon signal, for example for a pure Higgsino or wino LSP, or via direct detection of galactic dark matter, for example for the case of a mixed singlino/Higgsino LSP. The former case would motivate the construction of a lepton collider with energy optimized for the particular energy of the photon signal, while the latter would motivate very high luminosity LHC studies. A key goal would be to learn sufficient about the structure of the underlying theory to demonstrate that electroweak symmetry breaking has a very high degree of fine-tuning, pointing to a multiverse.

The simplest and most natural theories of weak-scale supersymmetry lead to a Higgs boson lighter than about 100 GeV; of course, theories can be extended and naturalness can be relaxed to accommodate a heavier Higgs. On the contrary, the theories introduced in this paper *predict* a heavier Higgs in the mass range $\approx 120 - 145$, $110 - 140$, and $100 - 130$ GeV for Higgsino LSP Spread Supersymmetry, Wino LSP Spread Supersymmetry, and the Environmental MSSM, respectively. We expect this Higgs boson to be indistinguishable from the Standard Model Higgs, and for colored superpartners to be out of reach for the LHC.

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